

# Pathways to testable leptogenesis

Pei-Hong Gu<sup>1\*</sup> and Utpal Sarkar<sup>2†</sup>

<sup>1</sup>*The Abdus Salam International Centre for Theoretical Physics, Strada Costiera 11, 34014 Trieste, Italy*

<sup>2</sup>*Physical Research Laboratory, Ahmedabad 380009, India*

In the conventional seesaw models of neutrino masses, leptogenesis occurs at a very high scale. Three approaches have been discussed in the literature to lower the scale of leptogenesis: mass degeneracy, hierarchy of couplings and three-body decays. We advocate yet another approach to a testable leptogenesis, whereby the decaying particles could go out of equilibrium at an accessible scale due to kinematics, although their couplings to the decay products are larger for generating a desired CP asymmetry. We demonstrate this new possibility for the testable leptogenesis in a two Higgs doublet model where the neutrino masses originate from a one-loop diagram.

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Among various baryogenesis mechanisms, which solve the puzzle of the matter-antimatter asymmetry in the universe, the electroweak baryogenesis is the most attractive one because of its testability. However, it died in the standard model (SM) and only survives in the minimal supersymmetric model (MSSM) with some rather restrictive conditions or in the context of more complicated models.

At present, the leptogenesis [1, 2, 3, 4, 5, 6, 7, 8, 9] scenario, where a lepton asymmetry is firstly generated and then is partially converted to a baryon asymmetry by the sphalerons [10], is definitely the simplest baryogenesis theory to test in the near future. Usually, the leptogenesis models also explain the tiny but nonzero neutrino masses by the seesaw mechanism [11, 12, 13, 14, 15, 16]. Since the tiny neutrino mass is naturally generated by a large lepton number violating scale in the seesaw models, the decays of heavier particles at the large lepton number violating scale also generates the leptogenesis in these models.

Attempts have been made to lower the scale of leptogenesis so as to make them testable in the next generation accelerators. Three pathways to a testable leptogenesis model have been demonstrated in some variants of the seesaw models through out-of-equilibrium CP-violating decays of some TeV scale particles, which could be detected at the LHC or ILC. These possibilities are: mass degeneracy [5, 7] (the decaying particles are assumed to be almost degenerate), hierarchy of couplings [17, 18] and three-body decays [19]. In the case of mass degeneracy, the decaying particles have tiny couplings to guarantee their decay widths smaller than the expansion rate of the universe, parametrized in terms of the Hubble constant to satisfy the out-of-equilibrium condition. The decaying particles have quasi-degenerate masses to resonantly enhance the CP asymmetry. In the case of hierarchical couplings, the lighter decaying and heavier virtual particles have smaller and larger couplings, respectively, thus the

smaller couplings determine the decay widths while the large couplings dominate the CP asymmetry. In the case of three-body decays, the decay widths of the lighter particles are suppressed by the heavy masses of the virtual particles, whose couplings are large enough to enhance the CP asymmetry. Note the decaying particles should be gauge singlets of any low-energy gauge symmetry, otherwise, the scattering processes originating from the gauge interactions will considerably damp the produced lepton asymmetry.

Obviously, the three pathways have a same essence that the (effective) couplings of the decaying particles to the decay products should be much smaller than unity so that the decay widths could be smaller than the Hubble constant at a low scale. This conclusion is based on an assumption that the decay products are much lighter than the decaying particles and hence the decay widths are only related to the masses of the decaying particles and the couplings. However, it is not the only possibility because the decay widths could be kinematically suppressed even if the couplings are larger.

In this paper, we discuss the kinematical effect in the leptogenesis scenario, which could alter the out-of-equilibrium condition considerably. For demonstration, we consider a two Higgs doublet model [20], where the neutrino masses originate from a one-loop diagram. Although the original model was proposed to include a dark matter candidate, we relax this condition to allow a different range of the parameters. With the larger couplings and the proper masses of the decaying particles and decay products, the sizable CP asymmetry and the suppressed decay widths are simultaneously allowed to realize a successful leptogenesis at an accessible scale.

We now introduce the quoted two Higgs doublet model [20], where the SM is extended by three right-handed neutrinos  $N_R$  and one scalar doublet  $\eta = (\eta^0, \eta^-)^T$ . The new particles are odd under a  $Z_2$  discrete symmetry. As a result, the following Yukawa couplings and Majorana mass term are allowed,

$$\mathcal{L} \supset -y\overline{\psi}_L\eta N_R - \frac{1}{2}M_N\overline{N}_R^c N_R + \text{h.c.}, \quad (1)$$

where  $\psi_L = (\nu_L, \ell_L)^T$  denotes the SM left-handed lep-

\*Electronic address: pgu@ictp.it

†Electronic address: utpal@prl.res.in

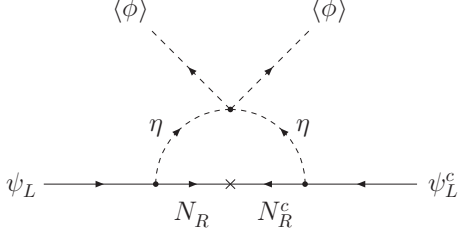


FIG. 1: The one-loop diagram for generating the radiative neutrino masses.

tons. For convenience and without loss of any generality, we will choose the basis in which the Majorana mass Matrix  $M_N$  is real and diagonal.

There are no Dirac masses between the left- and right-handed neutrinos since the exact  $Z_2$  protects  $\eta$  against any vacuum expectation values (*vevs*). So, the neutrinos will remain massless at tree level unless the quartic scalar interaction,

$$\mathcal{L} \supset -\frac{1}{2}\lambda \left[ (\phi^\dagger \eta)^2 + \text{h.c.} \right], \quad (2)$$

is introduced to realize the one-loop diagram as depicted in Fig. 1. Here  $\phi = (\phi^0, \phi^-)^T$  is the SM Higgs with  $\langle \phi \rangle \equiv v \simeq 174 \text{ GeV}$ . For the purpose of demonstration, we will take  $\lambda$  to be positive in the following discussions and calculations. The radiative neutrino masses can be explicitly calculated [20],

$$(m_\nu)_{ij} = \frac{1}{16\pi^2} \sum_{k=1}^3 y_{ik} y_{jk} M_{N_k} \times \left[ \frac{M_{\eta_R^0}^2}{M_{\eta_R^0}^2 - M_{N_k}^2} \ln \left( \frac{M_{\eta_R^0}^2}{M_{N_k}^2} \right) - \frac{M_{\eta_I^0}^2}{M_{\eta_I^0}^2 - M_{N_k}^2} \ln \left( \frac{M_{\eta_I^0}^2}{M_{N_k}^2} \right) \right], \quad (3)$$

where  $\eta_R^0$  and  $\eta_I^0$  are defined by  $\eta^0 = \frac{1}{\sqrt{2}} (\eta_R^0 + i\eta_I^0)$ . For  $M_{\eta_R^0}^2 - M_{\eta_I^0}^2 = 2\lambda v^2 \ll M_{\eta_{R,I}^0}^2$  and  $M_{N_k}^2 > M_{\eta_{R,I}^0}^2$ , we can obtain a simple formula for the neutrino masses,

$$(m_\nu)_{ij} = \frac{\mathcal{O}(\lambda)}{16\pi^2} \sum_{k=1}^3 \frac{y_{ik} y_{jk}}{M_{N_k}} v^2. \quad (4)$$

Therefore, the observed neutrino masses can be naturally explained, for example,  $m_\nu \sim \mathcal{O}(0.01 - 0.1 \text{ eV})$  for  $\lambda = \mathcal{O}(10^{-4})$ ,  $y \sim \mathcal{O}(10^{-3})$  and  $M_{N_i} = \mathcal{O}(100 \text{ GeV} - 10 \text{ TeV})$ .

The heavy Majorana neutrinos  $N = N_R + N_R^c$  can decay into the SM left-handed leptons and the doublet scalar. For simplicity, only the decay channel of  $N \rightarrow \psi_L \eta^*$  is shown in Fig. 2 and its CP conjugate state  $N \rightarrow$

$\psi_L^c \eta$  is omitted. At the tree level, the decay widths are given by

$$\Gamma(N_i \rightarrow \psi_L + \eta^*) = \Gamma(N_i \rightarrow \psi_L^c + \eta) \quad (5)$$

$$= \frac{1}{16\pi} (y^\dagger y)_{ii} M_{N_i} r_{N_i}^2, \quad (6)$$

where the factor,

$$r_{N_i} = 1 - \frac{M_\eta^2}{M_{N_i}^2}, \quad (7)$$

depends on the mass of the decay product in addition to that of the decaying particle. For the increasing values of  $M_\eta$ , we get a smooth interpolation from  $r_{N_i} \simeq 1$  ( $M_\eta^2 \ll M_{N_i}^2$ ) to  $r_{N_i} \rightarrow 0$  ( $M_\eta^2 \rightarrow M_{N_i}^2$ ). Usually the decay products are assumed to be very light, and hence, this factor is taken to be 1. However, this factor can also be very small. We shall consider the case when this factor is very small. In the present model under consideration, we shall not provide any explanation for the smallness of this parameter  $r_{N_i}$ . But for purpose of demonstration let us give an explicit example, where this parameter could be very small naturally.

Consider a supersymmetric model with superfields  $N$  and  $S$ . An Yukawa coupling  $NNS$  would then allow the decay of the fermionic component of  $N$  to its lighter scalar component and the fermionic component of  $S$ :  $N_f \rightarrow N_s + S_f$ . If the masses of these fields are  $M_N \sim 10^7 \text{ GeV}$  and  $M_S \ll M_N$ , the mass splitting between  $N_f$  and  $N_s$  will be of the order of supersymmetry breaking scale of  $10^3 \text{ GeV}$ , so that this decay width will have a suppression factor of  $r_N \sim 10^{-4}$ . Although this example can not be extrapolated to the present model, we shall assume a similar small value for  $r_{N_i}$ .

Coming back to the present model, the decays of the heavy Majorana neutrinos can create a lepton asymmetry, if the Yukawa couplings  $y$  provide the source of the CP violation. The lowest order non-trivial asymmetry comes from the interference of the tree-level diagrams with the one-loop diagrams. We calculate the CP asymmetry and find it to be same as that in the canonical seesaw model,

$$\begin{aligned} \varepsilon_{N_i} &\equiv \frac{\Gamma(N_i \rightarrow \psi_L + \eta^*) - \Gamma(N_i \rightarrow \psi_L^c + \eta)}{\Gamma(N_i \rightarrow \psi_L + \eta^*) + \Gamma(N_i \rightarrow \psi_L^c + \eta)} \\ &\simeq \frac{1}{8\pi} \frac{1}{(y^\dagger y)_{ii}} \sum_{j \neq i} \text{Im} \left[ (y^\dagger y)_{ij}^2 \right] \\ &\quad \times \left[ f \left( \frac{M_{N_j}^2}{M_{N_i}^2} \right) + g \left( \frac{M_{N_j}^2}{M_{N_i}^2} \right) \right], \end{aligned} \quad (8)$$

which is free from the masses of the decay products. Here the functions  $f$  and  $g$  are the contributions from the vertex and self-energy corrections, respectively:

$$f(x) = \sqrt{x} \left[ 1 - (1+x) \ln \left( \frac{1+x}{x} \right) \right], \quad (9)$$

$$g(x) = \frac{\sqrt{x}}{1-x}. \quad (10)$$

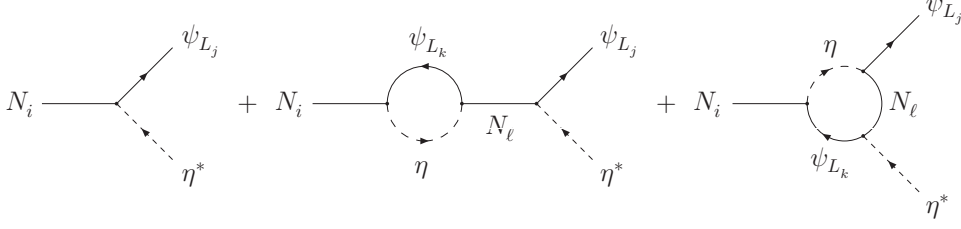


FIG. 2: The heavy Majorana neutrinos decay into the SM left-handed leptons and the doublet scalar. Their CP conjugate states have been omitted for simplicity.

For illustration, we consider the limiting case with  $M_{N_1} \ll M_{N_{2,3}}$  (a factor  $M_{N_{2,3}}/M_{N_1}$  of 3 – 10 is enough because the number density of  $N_{2,3}$  is fastly Boltzmann suppressed at the temperature below its mass), where the final lepton asymmetry should mainly come from the contributions of the decays of  $N_1$ . We can simplify the CP asymmetry (8) as

$$\varepsilon_{N_1} \simeq -\frac{3}{16\pi} \frac{1}{(y^\dagger y)_{11}} \sum_{j=2,3} \text{Im} \left[ (y^\dagger y)_{1j}^2 \right] \frac{M_{N_1}}{M_{N_j}}. \quad (11)$$

Similar to the DI bound [21, 22], one can also deduce an upper bound on the above CP asymmetry by inserting the neutrino masses (4) into the CP asymmetry (11),

$$|\varepsilon_{N_1}| < \frac{3\pi}{\mathcal{O}(\lambda)} \frac{M_{N_1} m_3}{v^2} |\sin \delta| \quad (12)$$

with  $m_3$  and  $\delta$  being the biggest eigenvalue of the neutrino mass matrix and the CP phase, respectively. Here we have assumed the neutrinos to be hierarchical [23]. Clearly, in the present case, the DI bound is relaxed by a factor of  $16\pi^2/\mathcal{O}(\lambda)$ .

For effectively creating a lepton asymmetry in the thermal evolution of the universe, the decays of  $N_1$  should satisfy the condition of departure from equilibrium, which is described by

$$\Gamma_{N_1} \lesssim H(T) \Big|_{T=M_{N_1}}. \quad (13)$$

where

$$\begin{aligned} \Gamma_{N_1} &= \Gamma(N_1 \rightarrow \psi_L + \eta^*) + \Gamma(N_1 \rightarrow \psi_L^c + \eta) \\ &= \frac{1}{8\pi} (y^\dagger y)_{11} M_{N_1} r_{N_1}^2 \end{aligned} \quad (14)$$

is the total decay width of  $N_1$  and

$$H(T) = \left( \frac{8\pi^3 g_*}{90} \right)^{\frac{1}{2}} \frac{T^2}{M_{\text{Pl}}} \quad (15)$$

is the Hubble constant with the Planck mass  $M_{\text{Pl}} \simeq 10^{19}$  GeV and the relativistic degrees of freedom  $g_* \simeq 100$  [24].

It is straightforward to perform the out-of-equilibrium condition by inserting (14) and (15) to (13),

$$(y^\dagger y)_{11} \lesssim \left( \frac{256 \cdot \pi^5 \cdot g_*}{45} \right)^{\frac{1}{2}} \frac{M_{N_1}}{M_{\text{Pl}}} \frac{1}{r_{N_1}^2} \quad (16)$$

for  $y \sim \mathcal{O}(10^{-3})$ ,  $M_{N_1} = \mathcal{O}(100 \text{ GeV} - 1 \text{ TeV})$  and  $r_{N_1} = \mathcal{O}(10^{-4})$ . So, by inputting  $\lambda = 10^{-4}$ ,  $M_{N_1} = 700 \text{ GeV}$ ,  $m_3 = 0.07 \text{ eV}$  and  $\sin \delta = -1$  to the CP asymmetry (12), we derive  $\varepsilon_{N_1} \simeq -1.5 \times 10^{-7}$  and then obtain the final baryon asymmetry,

$$\begin{aligned} \frac{n_B}{s} &= \frac{28}{79} \frac{n_{B-L}}{s} = -\frac{28}{79} \frac{n_L}{s} \\ &\simeq -\frac{28}{79} \varepsilon_{N_1} \frac{n_{N_1}^{eq}}{s} \Big|_{T=M_{N_1}} \simeq -\frac{1}{15} \frac{\varepsilon_{N_1}}{g_*} \\ &\simeq 10^{-10} \end{aligned} \quad (17)$$

as desired to explain the matter-antimatter asymmetry of the universe.

In this leptogenesis scenario, all of the new particles are at an accessible scale and then are expected to be detected at the forthcoming experiments. Specifically, the  $\eta$  scalars can be directly produced in pairs by the SM gauge bosons  $W^\pm$ ,  $Z$  or  $\gamma$ . Once produced,  $\eta^\pm$  will decay into  $\eta_{R,I}^0$  and a virtual  $W^\pm$ , which becomes a quark-antiquark or lepton-antilepton pair. If  $\eta_R^0$  is heavier than  $\eta_I^0$  for  $\lambda > 0$ , the decay chain,

$$\eta^+ \rightarrow \eta_R^0 l^+ \nu, \quad \text{then} \quad \eta_R^0 \rightarrow \eta_I^0 l^+ l^-, \quad (18)$$

has 3 charged leptons and large missing energy, and can be compared to the direct decay,

$$\eta^+ \rightarrow \eta_I^0 l^+ \nu, \quad (19)$$

to extract the masses of the respective particles.

As for the heavy Majorana neutrinos  $N_{1,2,3}$ , they can be produced in pairs by the  $e^+e^-$  annihilation through the  $\eta^\pm$  exchanges as shown in Fig. 3, which could be significant for the larger Yukawa couplings. We give the cross sections as below,

$$\sigma(e^+e^- \rightarrow N_i N_i) = \frac{1}{32\pi} |y_{ei}|^4 (F_{ii} + H_{ii}), \quad (20)$$

$$\sigma(e^+e^- \rightarrow N_i N_j) = \frac{1}{32\pi} |y_{ei}|^2 |y_{ej}|^2 F_{ij}, \quad (21)$$

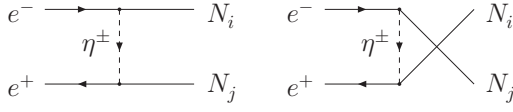


FIG. 3: The heavy Majorana neutrinos  $N_{1,2,3}$  are produced in pairs by the  $e^+e^-$  annihilation through the  $\eta^\pm$  exchanges,

where

$$F_{ij} = \frac{1}{s(s-4m_e^2)} \left\{ \left( 2M_{\eta^\pm}^2 - 2m_e^2 - M_{N_i}^2 - M_{N_j}^2 \right) \times \ln \left| \frac{t_0 - M_{\eta^\pm}^2}{t_1 - M_{\eta^\pm}^2} \right| + \left[ \frac{(M_{\eta^\pm}^2 - M_{N_i}^2 - m_e^2)(M_{\eta^\pm}^2 - M_{N_j}^2 - m_e^2)}{(t_0 - M_{\eta^\pm}^2)(t_1 - M_{\eta^\pm}^2)} + 1 \right] (t_0 - t_1) \right\}, \quad (22)$$

$$(23)$$

$$H_{ii} = \frac{M_{N_i}^2 (s - 2m_e^2)}{s(s-4m_e^2)(s-2m_e^2-2M_{N_i}^2+2M_{\eta^\pm}^2)} \times \ln \left| \frac{t_0 - M_{\eta^\pm}^2}{t_1 - M_{\eta^\pm}^2} \right|, \quad (24)$$

with

$$t_0(t_1) = -\frac{1}{4} \left\{ (s-4m_e^2)^{\frac{1}{2}} \mp \left[ \frac{s - (M_{N_i} - M_{N_j})^2}{s} \right]^{\frac{1}{2}} \times \left[ s - (M_{N_i} + M_{N_j})^2 \right]^{\frac{1}{2}} \right\}^2. \quad (25)$$

In this paper, we point out the kinematical approach to the testable leptogenesis. The kinematical effect allows the (effective) couplings of the decaying particles to the decay products to be larger for the sizable CP asymmetry, meanwhile, allows the decay widths to be smaller for the out-of-equilibrium condition at an accessible scale. We demonstrate this scenario in a two Higgs doublet model where the Majorana neutrino masses originate from a one-loop diagram. We can also apply the kinematical effect to the model where the neutrinos are Dirac particles and their tiny masses are radiatively generated [25]. These leptogenesis models can be verified in the near future.

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